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Cosmological Models in the Nonsymmetric Gravitational Theory

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Abstract

Cosmological consequences of the nonsymmetric gravitational theory (NGT) are studied. The structure of the NGT field equations is analyzed for an inhomogeneous and anisotropic universe, based on the spherically symmetric field equations. It is assumed that the matter density and pressure are purely time dependent, and it is shown that the field equations allow open, flat and closed universes. The field equations are expanded about the Friedmann-Robertson-Walker (FRW) model for a small antisymmetric field. The observations are most in accord with the approximately spatially flat universe, which has an FRW metric to lowest order in the skew field. The recent analysis of galaxy synchroton radiation polarization data, which indicates a residual polarization after Faraday polarization is subtracted, can be explained by the birefringence of electromagnetic waves produced by the NGT spacetime.

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I. INTRODUCTION

A new version of the nonsymmetric gravitational theory (NGT) has been published [1–10], which possesses a linear approximation free of ghost poles and tachyons, with well-behaved asymptotic conditions and a stable expansion about a general relativistic (GR) background to first order in the skew field $g_{[\mu\nu]}$.

The standard model scenario in big bang cosmology assumes that at any given time the universe is homogeneous and isotropic when averaged over a sufficiently large scale. The equations of the standard model can be expressed as

$$H^2 + \frac{k}{R^2} = \Omega_M H^2 + \Omega_\Lambda H^2, \tag{1}$$

where $H = \dot{R}/R$ is the Hubble parameter, Ω_M is the density parameter for matter and Ω_Λ is the density parameter contributed by the cosmological constant. This model gives a good account of present experimental data [12], although there are indications that the latest data for the age of the universe may require a non-zero cosmological constant [13].

In the following, we shall derive a cosmological model in which it is assumed that the skew $g_{[\mu\nu]}$ contributions are small in the present universe, so that we can expand the NGT field equations about the homogeneous and isotropic Friedmann-Robertson-Walker (FRW) model and obtain an approximate dynamical equation for the cosmological scale factor $R(t)$. We base the cosmological solutions on an inhomogeneous spherically symmetric $g_{\mu\nu}$, and assume for simplicity that the matter density ρ_M and the pressure p are uniform and only depend on the time t. For the case of a spatially flat universe, the metric takes the simple form of an FRW universe to lowest order in $g_{[\mu\nu]}$ and the Friedmann-type equation for $R(t)$ contains corrections due to a non-zero $g_{[\mu\nu]}$ and its derivatives.

Savaria [11] obtained an exact plane symmetric anisotropic solution both for the vacuum field equations and for the field equations in the presence of pressureless matter.

II. STRUCTURE OF THE NGT FIELD EQUATIONS

We shall decompose the nonsymmetric $g_{\mu\nu}$ and $\Gamma^\lambda_{\mu\nu}$ as

$$g_{(\mu\nu)} = \frac{1}{2}(g_{\mu\nu} + g_{\nu\mu}), \quad g_{[\mu\nu]} = \frac{1}{2}(g_{\mu\nu} - g_{\nu\mu}), \tag{2}$$

and

$$\Gamma^\lambda_{\mu\nu} = \Gamma^\lambda_{(\mu\nu)} + \Gamma^\lambda_{[\mu\nu]}. \tag{3}$$

The contravariant tensor $g^{\mu\nu}$ is defined in terms of the equation:

$$g^{\mu\nu}g_{\sigma\nu} = g^{\nu\mu}g_{\nu\sigma} = \delta^\mu{}_\sigma. \tag{4}$$

We shall use the notation: $\mathbf{g}^{\mu\nu} = \sqrt{-g}g^{\mu\nu}$, $g = \mathrm{Det}(g_{\mu\nu})$, and $R_{\mu\nu}(W)$ is the NGT contracted curvature tensor:

$$R_{\mu\nu}(W) = W^\beta_{\mu\nu,\beta} - \frac{1}{2}(W^\beta_{\mu\beta,\nu} + W^\beta_{\nu\beta,\mu}) - W^\beta_{\alpha\nu}W^\alpha_{\mu\beta} + W^\beta_{\alpha\beta}W^\alpha_{\mu\nu}, \tag{5}$$

defined in terms of the unconstrained nonsymmetric connection:

$$W^\lambda_{\mu\nu} = \Gamma^\lambda_{\mu\nu} - \frac{2}{3}\delta^\lambda{}_\mu W_\nu, \tag{6}$$

where

$$W_\mu = \frac{1}{2}(W^\lambda_{\mu\lambda} - W^\lambda_{\lambda\mu}).$$

Eq.(6) leads to the result:

$$\Gamma_\mu = \Gamma^\lambda_{[\mu\lambda]} = 0.$$

The contracted tensor $R_{\mu\nu}(W)$ can be written as

$$R_{\mu\nu}(W) = R_{\mu\nu}(\Gamma) + \frac{2}{3}W_{[\mu,\nu]},$$

where

$$R_{\mu\nu}(\Gamma) = \Gamma^{\beta}_{\mu\nu,\beta} - \frac{1}{2}\left(\Gamma^{\beta}_{(\mu\beta),\nu} + \Gamma^{\beta}_{(\nu\beta),\mu}\right) - \Gamma^{\beta}_{\alpha\nu}\Gamma^{\alpha}_{\mu\beta} + \Gamma^{\beta}_{(\alpha\beta)}\Gamma^{\alpha}_{\mu\nu}.$$

The field equations in the new version of NGT take the form [1]:

$$G_{\mu\nu}(W) + \Lambda g_{\mu\nu} + S_{\mu\nu} = 8\pi(T_{\mu\nu} + K_{\mu\nu}), \tag{7a}$$

$$\mathbf{g}^{[\mu\nu]}{}_{,\nu} = -\frac{1}{2}\mathbf{g}^{(\mu\alpha)}W_{\alpha}, \tag{7b}$$

$$\mathbf{g}^{[\mu\nu]}u_{\nu} = 0, \tag{7c}$$

$$\begin{aligned}\mathbf{g}^{\mu\nu}{}_{,\sigma} + \mathbf{g}^{\rho\nu}W^{\mu}_{\rho\sigma} + \mathbf{g}^{\mu\rho}W^{\nu}_{\sigma\rho} - \mathbf{g}^{\mu\nu}W^{\rho}_{\sigma\rho} + \frac{2}{3}\delta^{\nu}_{\sigma}\mathbf{g}^{\mu\rho}W^{\beta}_{[\rho\beta]} \\ +\frac{1}{6}(\mathbf{g}^{(\mu\beta)}W_{\beta}\delta^{\nu}_{\sigma} - \mathbf{g}^{(\nu\beta)}W_{\beta}\delta^{\mu}_{\sigma}) = 0.\end{aligned} \tag{7d}$$

Here, we have

$$G_{\mu\nu} = R_{\mu\nu} - \frac{1}{2}g_{\mu\nu}R, \tag{8a}$$

$$S_{\mu\nu} = -\frac{1}{6}(W_{\mu}W_{\nu} - \frac{1}{2}g_{\mu\nu}g^{\alpha\beta}W_{\alpha}W_{\beta}), \tag{8b}$$

$T_{\mu\nu}$ denotes the nonsymmetric source tensor and Λ is the cosmological constant. Moreover, using Eq.(7c) we have

$$K_{\mu\nu} = -\frac{1}{8\pi}u_{[\mu}\phi_{\nu]}, \tag{9}$$

where ϕ_{μ} are four Lagrange multiplier fields [1], $u_{\mu} = dx_{\mu}/d\tau$ is the four-velocity vector and τ denotes the proper time along an observer's world line in spacetime. We impose the condition:

$$g^{(\mu\nu)}u_{\mu}u_{\nu} = 1. \tag{10}$$

We can choose the vector u_{μ} to be $u_{\mu} = (0, 0, 0, 1/\sqrt{g^{00}})$, so that (7c) corresponds to the three constraint equations:

$$\mathbf{g}^{[i0]} = 0. \tag{11}$$

The generalized Bianchi identities:

$$[\mathbf{g}^{\alpha\nu}G_{\rho\nu}(\Gamma) + \mathbf{g}^{\nu\alpha}G_{\nu\rho}(\Gamma)]_{,\alpha} + g^{\mu\nu}{}_{,\rho}\mathbf{G}_{\mu\nu} = 0, \tag{12}$$

give rise to the matter response equations [4]:

$$g_{\mu\rho}\mathbf{T}^{\mu\nu}{}_{,\nu} + g_{\rho\mu}\mathbf{T}^{\nu\mu}{}_{,\nu} + (g_{\mu\rho,\nu} + g_{\rho\nu,\mu} - g_{\mu\nu,\rho})\mathbf{T}^{\mu\nu} = 0. \tag{13}$$

After eliminating the Lagrange multiplier ϕ_μ from the field equations (7a), we get

$$G_{(\mu\nu)}(W) + \Lambda g_{(\mu\nu)} + S_{(\mu\nu)} = 8\pi T_{(\mu\nu)}, \tag{14a}$$

$$\epsilon^{\mu\nu\alpha\beta}u_\alpha(G_{[\mu\nu]} + \Lambda g_{[\mu\nu]} + S_{[\mu\nu]}) = 8\pi\epsilon^{\mu\nu\alpha\beta}u_\alpha T_{[\mu\nu]}, \tag{14b}$$

where $\epsilon^{\mu\nu\alpha\beta}$ is the Levi-Civita symbol.

For the case of a spherically symmetric field, the canonical form of $g_{\mu\nu}$, in NGT, is given by

$$g_{\mu\nu} = \begin{pmatrix} -\alpha & 0 & 0 & w \\ 0 & -\beta & f\sin\theta & 0 \\ 0 & -f\sin\theta & -\beta\sin^2\theta & 0 \\ -w & 0 & 0 & \gamma \end{pmatrix},$$

where α, β, γ and w are functions of r and t. From the constraint equations (11), we have $w = 0$ and only the $g_{[23]}$ component of $g_{[\mu\nu]}$ is different from zero. We have

$$\sqrt{-g} = \sin\theta[\alpha\gamma(\beta^2 + f^2)]^{1/2}.$$

If we adopt the approximation scheme leading to the geodesic equation for falling test particles, then we can use a comoving coordinate system with the velocity components [6]:

$$u^0 = 1, \quad u^r = u^\theta = u^\phi = 0,$$

and the time dependent metric in normal Gaussian form:

$$ds^2 = dt^2 - \alpha(r,t)dr^2 - \beta(r,t)(d\theta^2 + \sin^2\theta d\phi^2). \tag{15}$$

The field equations for the spherically symmetric system take the form:

$$G_{\mu\nu}(\Gamma) + \Lambda g_{\mu\nu} = 8\pi G T_{\mu\nu}, \tag{16}$$

$$\mathbf{g}^{[\mu\nu]}{}_{,\nu} = 0, \tag{17}$$

We have for the spherically symmetric conservation laws:

$$\mathrm{Re}[(T_{\rho-\nu-}\mathbf{g}^{\sigma+\nu-})_{;\sigma}] = 0, \tag{18}$$

where we have used the Einstein + and - notation for covariant differentiation with respect to the $\Gamma^\lambda_{\mu\nu}$ connection [14]. We assume the energy-momentum tensor takes the form:

$$T^{\mu\nu} = (\rho + p)u^\mu u^\nu - pg^{\mu\nu} + N^{[\mu\nu]}, \tag{19}$$

where $N^{[\mu\nu]}$ is a direct source for the antisymmetric fields.

In order to simplify the field equations, we must make several approximations. We shall assume that the direct coupling term in the source tensor, $N^{[\mu\nu]}$, is small and can be neglected and we also set $\Lambda = 0$.

To facilitate a comparison of the NGT cosmological models with the standard FRW model, we shall adopt the approximation:

$$\beta(r,t) \gg f(r,t). \tag{20}$$

We obtain the field equations:

$$-\frac{1}{\alpha}\left[\frac{\beta''}{\beta} - \frac{\beta'^2}{2\beta^2} - \frac{\alpha'\beta'}{2\alpha\beta}\right] + \frac{\ddot{\alpha}}{2\alpha} - \frac{\dot{\alpha}^2}{4\alpha^2} + \frac{\dot{\alpha}\dot{\beta}}{2\alpha\beta} + W = 4\pi G(\rho_M - p), \tag{21a}$$

$$\frac{1}{\beta} - \frac{1}{\alpha}\left(\frac{\beta''}{2\beta} - \frac{\alpha'\beta'}{4\alpha\beta}\right) + \frac{\ddot{\beta}}{2\beta} + \frac{\dot{\alpha}\dot{\beta}}{4\alpha\beta} + X = 4\pi G(\rho_M - p), \tag{21b}$$

$$-\frac{\ddot{\alpha}}{2\alpha} - \frac{\ddot{\beta}}{\beta} + \frac{\dot{\alpha}^2}{4\alpha^2} + \frac{\dot{\beta}^2}{2\beta^2} + Y = 4\pi G(\rho_M + 3p), \tag{21c}$$

$$-\frac{\dot{\beta}'}{\beta} + \frac{\beta'\dot{\beta}}{2\beta^2} + \frac{\dot{\alpha}\beta'}{2\alpha\beta} + Z = 0, \tag{21d}$$

$$\begin{aligned} &\frac{\dot{f}\dot{\beta}}{2\beta f} - \frac{\ddot{f}}{2f} + \frac{\beta'^2}{2\alpha\beta^2} - \frac{\dot{\beta}^2}{2\beta^2} + \frac{\alpha'\beta'}{2\alpha^2\beta} - \frac{\beta''}{\alpha\beta} + \frac{\ddot{\beta}}{\beta} - \frac{f'\beta'}{2\alpha\beta f} \\ &\qquad - \frac{\dot{\alpha}\dot{f}}{4\alpha f} + \frac{\dot{\alpha}\dot{\beta}}{2\alpha\beta} + \frac{f''}{2\alpha f} - \frac{\alpha' f'}{4\alpha^2 f} = 4\pi G(\rho_M - p). \end{aligned} \tag{21e}$$

Here, we have defined $\dot{\alpha} = d\alpha/dt$, $\alpha' = d\alpha/dr$ and

$$W(r,t) = -\frac{\alpha'\beta' f^2}{2\alpha^2\beta^3} + \frac{\beta'' f^2}{\alpha\beta^3}$$

$$-\frac{\dot{\alpha}\dot{\beta}f^2}{2\alpha\beta^3}-\frac{5\beta'^2f^2}{2\alpha\beta^4}+\frac{\dot{\alpha}f\dot{f}}{2\alpha\beta^2}+\frac{\alpha'ff'}{2\alpha^2\beta^2}-\frac{ff''}{\alpha\beta^2}+\frac{4ff'\beta'}{\alpha\beta^3}-\frac{3f'^2}{2\alpha\beta^2}, \tag{22a}$$

$$\begin{aligned} X(r,t) = &-\frac{\dot{\alpha}\dot{\beta}f^2}{2\alpha\beta^3}-\frac{\alpha'\beta'f^2}{2\alpha^2\beta^3}+\frac{\beta''f^2}{\alpha\beta^3}-\frac{\ddot{\beta}f^2}{\beta^3}+\frac{\dot{\alpha}f\dot{f}}{2\alpha\beta^2}+\frac{\dot{\beta}^2f^2}{\beta^4}-\frac{ff''}{\beta^2\alpha}-\frac{\beta'^2f^2}{\alpha\beta^4} \\ &+\frac{\alpha'ff'}{2\alpha^2\beta^2}+\frac{\dot{f}^2}{2\beta^2}+\frac{f\ddot{f}}{\beta^2}-\frac{f'^2}{2\alpha\beta^2}-\frac{3f\dot{f}\dot{\beta}}{2\beta^3}+\frac{3ff'\beta'}{2\alpha\beta^3}, \end{aligned} \tag{22b}$$

$$Y(r,t) = \frac{\ddot{\beta}f^2}{\beta^3}-\frac{5\dot{\beta}^2f^2}{2\beta^4}-\frac{3\dot{f}^2}{2\beta^2}+\frac{4\dot{\beta}f\dot{f}}{\beta^3}-\frac{f\ddot{f}}{\beta^2}, \tag{22c}$$

$$\begin{aligned} Z(r,t) = &\frac{\dot{\beta}'f^2}{\beta^3}-\frac{5\dot{\beta}\beta'f^2}{2\beta^4}-\frac{\dot{\alpha}\beta'f^2}{2\alpha\beta^3}+\frac{2\dot{\beta}ff'}{\beta^3}-\frac{f\dot{f}'}{\beta^2}-\frac{3f'\dot{f}}{2\beta^2} \\ &+\frac{\dot{\alpha}ff'}{2\alpha\beta^2}+\frac{2\beta'f\dot{f}}{\beta^3}. \end{aligned} \tag{22d}$$

III. ANALYSIS OF COSMOLOGICAL MODELS

An infinitesimal transformation of coordinates is described by

$$x'^{\mu} = x^{\mu} + \xi^{\mu}. \tag{23}$$

The coordinate transformation law for $g_{\mu\nu}$ is given by

$$g'_{\mu\nu}(x') = g_{\alpha\beta}(x)\frac{\partial x^{\alpha}}{\partial x'^{\mu}}\frac{\partial x^{\beta}}{\partial x'^{\nu}}, \tag{24}$$

and the condition of homogeneity is

$$g'_{\mu\nu}(x) = g_{\mu\nu}(x). \tag{25}$$

We now obtain the Killing equation:

$$g_{\mu\sigma}\xi^{\sigma}{}_{,\nu} + g_{\sigma\nu}\xi^{\sigma}{}_{,\mu} + g_{\mu\nu,\sigma}\xi^{\sigma} = 0. \tag{26}$$

The symmetric and skew parts of (26) are given by:

$$g_{(\mu\sigma)}\xi^{\sigma}{}_{,\nu} + g_{(\sigma\nu)}\xi^{\sigma}{}_{,\mu} + g_{(\mu\nu),\sigma}\xi^{\sigma} = 0, \tag{27}$$

and

$$g_{[\mu\sigma]}\xi^{\sigma}{}_{,\nu} + g_{[\sigma\nu]}\xi^{\sigma}{}_{,\mu} + g_{[\mu\nu],\sigma}\xi^{\sigma} = 0. \tag{28}$$

Eq.(28) gives

$$f\xi^{3}{}_{,1} = 0. \tag{29}$$

The Killing vectors take the form:

$$\xi^1 = q(r)[(a_1 \sin\phi + a_2 \cos\phi)\sin\theta + a_3\cos\theta], \tag{30a}$$

$$\xi^2 = \frac{q(r)}{r}[(a_1 \sin\phi + a_2 \cos\phi)\cos\theta - a_3\sin\theta], \tag{30b}$$

$$\xi^3 = \frac{q(r)}{r\sin\theta}(a_1\cos\phi - a_2\sin\phi), \tag{30c}$$

$$\xi^0 = 0, \tag{30d}$$

where $q(r)$ is a function that can be determined from the symmetric Killing equation. From the form of the Killing vectors, we see that the condition of homogeneity in NGT requires that

$$f(r,t) = 0. \tag{31}$$

It follows that all strictly homogeneous and isotropic solutions in NGT cosmology reduce to the FRW solutions of GR.

We shall simplify our model even further and assume that the density ρ_M and the pressure p are independent of position. It is assumed that a solution can be found by a separation of variables:

$$\alpha(r,t) = h(r)R^2(t), \quad \beta(r,t) = r^2S^2(t). \tag{32}$$

From Eq.(21d), we get

$$\frac{\dot{R}}{R} - \frac{\dot{S}}{S} = \frac{1}{2}Z(r,t)r. \tag{33}$$

If we assume that $Z(r,t) \approx 0$, then from (33) we find that

$$R(t) \approx S(t).$$

From the conservation law (18), we obtain within our approximation scheme:

$$\dot{p} = R^{-3}(t)\frac{\partial}{\partial t}[R^3(t)(\rho + p)].$$

Let us expand the metric $g_{(\mu\nu)}$ as

$$g_{(\mu\nu)}(r,t) = g^{HI}_{(\mu\nu)}(r,t) + \delta g_{(\mu\nu)}(r,t), \tag{34}$$

where $g^{HI}_{(\mu\nu)}$ denotes the homogeneous and isotropic solution of $g_{(\mu\nu)}$, and $\delta g_{(\mu\nu)}$ are small quantities which break the maximally symmetric solution, $g^{\rm HI}_{(\mu\nu)}$, to lowest order in $g_{[\mu\nu]}$. Within this scenario, the metric line-element takes the FRW form:

$$ds^2 = dt^2 - R^2(t)\Big[h(r)dr^2 + r^2(d\theta^2 + \sin^2\theta d\phi^2)\Big]. \tag{35}$$

Thus, we assume that the universe is approximately isotropic and homogeneous. We have three distinct solutions depending on whether $h(r) < 1, = 1, > 1$ and they correspond to hyperbolic (open), parabolic (flat), and elliptic (closed) universes, respectively.

Eqs.(21a) and (21c) can now be written:

$$2b(r) + \ddot{R}(t)R(t) + 2\dot{R}^2(t) + R^2(t)W(r,t) = 4\pi G R^2(t)(\rho_M(t) - p(t)), \tag{36}$$

$$-\ddot{R}(t)R(t) + \frac{1}{3}R^2(t)Y(t) = \frac{4\pi G}{3}R^2(t)(\rho_M(t) + 3p(t)), \tag{37}$$

where

$$b(r) = \frac{h'(r)}{2rh^2(r)}. \tag{38}$$

Eliminating $\ddot{R}$ by adding (36) and (37), we get

$$\dot{R}^2(t) = -b(r) + \frac{8\pi G}{3}[\rho_M(t) - \frac{3}{8\pi G}Q(r,t)]R^2(t), \tag{39}$$

where

$$Q(r,t) = \frac{1}{2}[W(r,t) + \frac{1}{3}Y(t)].$$

We can write Eq. (39) as

$$H^2(t) + \frac{b(r)}{R^2(t)} = \frac{8\pi G}{3}[\rho_M(t) - \frac{3}{8\pi G}Q(r,t)]. \tag{40}$$

This equation can in turn be written as

$$H^2(t) + \frac{b(r)}{R^2(t)} = \Omega(r,t)H^2(t),$$

where

$$\Omega(r,t) = \Omega_M(t) + \Omega_S(r,t),$$

and

$$\Omega_M(t) = \frac{8\pi G\rho_M(t)}{3H^2(t)}, \quad \Omega_S(r,t) = -\frac{Q(r,t)}{H^2(t)}.$$

Here, Ω_M denotes the familiar density parameter for matter, while Ω_S denotes the density parameter associated with the skew field contributions.

If $b(r) = 0$, then we get $\Omega(r,t) = 1$ and

$$H^2(t) = \frac{8\pi G}{3}[\rho_M(t) - \frac{3}{8\pi G}Q(t)]. \tag{41}$$

The line element now takes the approximate form of a flat, homogeneous and isotropic FRW universe:

$$ds^2 = dt^2 - R^2(t)[dr^2 + r^2(d\theta^2 + \sin^2\theta d\phi^2)]. \tag{42}$$

The field equations, Eq.(21), will determine $f(r,t) = g_{[23]}/\sin\theta$, which in turn will determine the function $b(r)$, given the matter density ρ_M. The present observations favour an approximately flat universe with $\Omega_0 \approx 1$.

A recent publication indicates that electromagnetic waves emitted in the form of synchrotron radiation from galaxies propagating over cosmological distances may suffer a residual polarization effect beyond the Faraday effect due to cosmic electromagnetic fields [15]. If upheld, these observations would have a possible explanation in the existence of an anisotropic effect due to the presence of the $g^{[23]}$ field component in the NGT-Maxwell action [16–20]:

$$I_{\mathrm{em}} = -\frac{1}{16\pi}\int d^4x\sqrt{-g}g^{\mu\alpha}g^{\nu\beta}F_{\mu\nu}F_{\alpha\beta}, \tag{43}$$

where $F_{\mu\nu}$ is the electromagnetic field strength.

The speed of light is not the same in every local inertial frame, breaking the strong equivalence principle in GR. Thus, the NGT spacetime acts like a crystal that produces birefringence as electromagnetic radiation passes through it. The residual rotation angle predicted by NGT is

$$\beta_S \approx \frac{1}{2}\Lambda_S(r,t)\cos^2(\gamma), \tag{44}$$

where γ is the angle between the propagating wave vector and the anisotropic axis fit to the data. A derivation and some consequences of the predicted polarization of electromagnetic waves due to the cosmological NGT spacetime will be discussed in a separate article [21].

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